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# STRONG INTERACTIONS OF WEAK BOSONS [a]

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We discuss the parameterization for electroweak gauge boson interactions without a light Higgs boson. We present the constraints on the anomalous gauge-boson couplings from the current experiments. We emphasize that the four-point couplings involving the longitudinal weak bosons are genuine to the underlying strong dynamics responsible for the electroweak symmetry breaking. We study the sensitivity to the four-point couplings and the possible heavy resonant states in this sector at future TeV $e^+e^-$ linear colliders.

## 1 Introduction

The most prominent mystery in contemporary particle physics is the origin of electroweak symmetry breaking (EWSB) and the mass generation mechanism for fermions. The celebrated Standard model (SM), which has passed the experimental test with high precision up to the energy scale $\mathcal{O}(100\ \text{GeV})$, fulfills the job by introducing an effective potential

$$V(\Phi) = \lambda(|\Phi|^2 - v^2/2)^2. \tag{1}$$

The scalar Higgs doublet can be parameterized by

$$\Phi = \exp(iw^a\tau^a/2v) \begin{pmatrix} 0 \\ (v+H)/\sqrt{2} \end{pmatrix},$$

where $w^a$ $(a = 1, 2, 3)$ are the Goldstone bosons and $H$ the physical Higgs boson of mass $m_H = \sqrt{2\lambda}v$. The vacuum expectation value (vev) $v \approx 246$ GeV sets the mass scale not only for the electroweak gauge bosons, but also for all fermions through Yukawa interactions. Thus searching for the Higgs boson has become the top priority to understand the electroweak symmetry breaking and mass generation mechanisms. However, due to the unknown parameter $\lambda$ in the Higgs potential, the Higgs boson mass $m_H$ is a free parameter. Current experimental searches at LEP2 have set a limit $m_H > 98.8$ GeV at a 95%

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confidence level.[1] Theoretically, the SM cannot be a consistent effective theory if $m_H > 800$ GeV. Thus, if there is no light Higgs boson found in the collider experiments, then new strong dynamics must set in. It has been pursued actively to explore the possible new dynamics responsible for the EWSB both theoretically and experimentally.

Although the original idea of Technicolor[2] for dynamical electroweak symmetry breaking is attractive, one would need the Extended Technicolor[3] to give fermions their masses. Tremendous efforts have been made to incorporate the vast fermion mass hierarchy yet to avoid severe flavor changing neutral currents.[4] Precision electroweak measurements at the $Z$ pole can also impose significant constraint on the new strong dynamics, such as on the number of Technicolors and the SU(2) breaking effects.[5] The fact that the top-quark mass is miraculously close to the electroweak scale makes it very attempting to consider the role of the top quark in the EWSB,[6] and the new interaction of topcolor has been also introduced to account for the EWSB and the top-quark mass generation.[7,8] More recently, models of dynamical symmetry breaking with seesaw mechanism of quark condensation are proposed.[9] Models of dynamical symmetry breaking often lead to predictions of rich phenomenology. Technicolor theories generically result in technihadrons like $\pi_T, \eta_T, \rho_T, A_{1T}$ and $\omega_T$. The topcolor models often have colorons ($Z'$, $V_8$); while the top seesaw models have flavorons ($\chi_{L,R}$, $F's$) and other composite scalars in the spectrum. If these particles exist well below 1 TeV, the experiments at the next generation of colliders will be able to discover them by their distinctive production and decay processes.

In this talk, we would like to take a different direction, namely a relatively model-independent approach to parameterize the Strongly-interacting Electro-Weak Sector (SEWS). One thing we know for sure in the gauge boson sector is that the longitudinally polarized states ($W_L^\pm, Z_L$) exist and they possess an (approximate) SU(2) custodial symmetry, leading to the mass relation $M_W \approx M_Z \cos\theta_W$ where $\theta_W$ is the weak mixing angle. In the scenario of spontaneous EWSB, the states $W_L^\pm, Z_L$ are equivalent to the Goldstone bosons $\omega^\pm, \omega^0$ at high energies.[10,11] One can thus construct an electroweak effective chiral Lagrangian based on the Goldstone bosons to parameterize the EWSB sector. One can also consider to include higher dimensional operators by the derivative expansion, which would lead to the "anomalous couplings" beyond the restricted form of the SM for the gauge bosons. We can also go a step further to introduce lower-lying heavy resonances near the TeV region. These will be discussed in the next section. In Sec. 3, we summarize the constraints on the anomalous coupling parameters from the current experimental results. We present analyses for the sensitivity to the SEWS sector at future

high energy/high luminosity $e^+e^-$ colliders in Sec. 4, where we emphasize the important role for the processes $W_L W_L \to W_L W_L$ as well as $W_L W_L \to t\bar{t}$. We make some concluding remarks in Sec. 5. For more discussions, there are recent reviews [12] that dealt with related topics.

## 2 Model-independent parameterization for SEWS

It is known that the most economical description of EWSB below a new physics scale is the electroweak chiral Lagrangian with non-linear realization of the symmetry.[13] The lowest order term respecting the symmetry can be written as

$$\mathcal{L}^{(2)} = \frac{v^2}{4}\mathrm{Tr}[(D^\mu U)^\dagger (D_\mu U)], \tag{2}$$

where

$$U = \exp[i\tau^a \omega^a / v], \qquad D_\mu U = \partial_\mu U + igW_\mu^a \frac{\tau^a}{2} U - ig' U B_\mu \frac{\tau^3}{2}. \tag{3}$$

$\mathcal{L}^{(2)}$ breaks the electroweak gauge symmetry spontaneously and the coefficient is fixed by the gauge-boson mass. It also respects the custodial SU(2) symmetry and the low-energy theorem [14] if gauge coupling $g'$ is ignored. To gain more information of the strong dynamics responsible for the EWSB, we need to go beyond the minimal term of $\mathcal{L}^{(2)}$. It should be mentioned that there is yet another popular approach to the electroweak effective Lagrangian, namely the linear realization with an explicit doublet Higgs field.[15] For the sake of simplicity, we will not discuss this approach. In the rest of this section, we will focus on some parameterization for the couplings among the Goldstone bosons, some lower-lying resonances as well as the top quark.

### *2.1 Scalar resonance*

One can parameterize an isosinglet heavy scalar resonance and its coupling to the Goldstone bosons and the top quark as

$$\mathcal{L}^H = \frac{1}{2}\partial^\mu H \partial_\mu H - \frac{1}{2}M_H^2 H^2 + \frac{1}{2}(g_h v H + g_h' \frac{H^2}{2})\mathrm{Tr}\partial^\mu U^\dagger \partial_\mu U - \kappa_t \frac{m_t}{v} H\bar{t}t. \tag{4}$$

There are four free parameters in this framework $M_H$, $g_h$, $g_h'$ and $\kappa_t$. The couplings $g_h$ and $\kappa_t$ can be traded to the partial decay widths to $W_L W_L$ and $t\bar{t}$, characterizing the dynamics of the underlying theory. This Lagrangian reduces to the SM Higgs sector when $g_h = g_h' = \kappa_t = 1$.

*2.2 Vector resonance*

There is a systematic way to introduce an isotriplet vector resonance into the chiral Lagrangian,[16] with an example called BESS [17] for the EWSB sector. There are three free parameters that can be expressed by the decay partial widths to $W_L W_L$ and to fermions, plus the vector boson mass $M_V$.[17,18] We note that other lower-lying vector states such as $A_1$ and $\omega_T$ can also be incorporated in the chiral Lagrangian.[19,20,21]

*2.3 Next-to-leading order terms*

If the resonant states are kinematically unaccessible, then the low-energy effects can be conveniently parameterized by the coefficients of higher order terms of the electroweak chiral Lagrangian as the derivative expansion. There are 12 $SU_L(2) \otimes U_Y(1)$ gauge invariant and CP-even operators,[13] parameterized by

$$
\begin{aligned}
\mathcal{L}^{(2)\prime} &= \ell_0(\frac{v}{\Lambda})^2 \ \frac{v^2}{4}[\mathrm{Tr}(\mathcal{T}\mathcal{V}_\mu)]^2 \ , \\
\mathcal{L}_1 &= \ell_1(\frac{v}{\Lambda})^2 \ \frac{gg'}{2} B_{\mu\nu}\mathrm{Tr}(\mathcal{T}\mathbf{W}^{\mu\nu}) \ , \\
\mathcal{L}_2 &= \ell_2(\frac{v}{\Lambda})^2 \ \frac{ig'}{2} B_{\mu\nu}\mathrm{Tr}(\mathcal{T}[\mathcal{V}^\mu, \mathcal{V}^\nu]) \ , \\
\mathcal{L}_3 &= \ell_3(\frac{v}{\Lambda})^2 \ ig\mathrm{Tr}(\mathbf{W}_{\mu\nu}[\mathcal{V}^\mu, \mathcal{V}^\nu]) \ , \\
\mathcal{L}_4 &= \ell_4(\frac{v}{\Lambda})^2[\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}_\nu)]^2 \ , \\
\mathcal{L}_5 &= \ell_5(\frac{v}{\Lambda})^2[\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}^\mu)]^2 \ , \\
\mathcal{L}_6 &= \ell_6(\frac{v}{\Lambda})^2[\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}_\nu)]\mathrm{Tr}(\mathcal{T}\mathcal{V}^\mu)\mathrm{Tr}(\mathcal{T}\mathcal{V}^\nu) \ , \\
\mathcal{L}_7 &= \ell_7(\frac{v}{\Lambda})^2[\mathrm{Tr}(\mathcal{V}_\mu\mathcal{V}^\mu)]\mathrm{Tr}(\mathcal{T}\mathcal{V}_\nu)\mathrm{Tr}(\mathcal{T}\mathcal{V}^\nu) \ , \\
\mathcal{L}_8 &= \ell_8(\frac{v}{\Lambda})^2 \ \frac{g^2}{4}[\mathrm{Tr}(\mathcal{T}\mathbf{W}_{\mu\nu})]^2 \ , \\
\mathcal{L}_9 &= \ell_9(\frac{v}{\Lambda})^2 \ \frac{ig}{2}\mathrm{Tr}(\mathcal{T}\mathbf{W}_{\mu\nu})\mathrm{Tr}(\mathcal{T}[\mathcal{V}^\mu, \mathcal{V}^\nu]) \ , \\
\mathcal{L}_{10} &= \ell_{10}(\frac{v}{\Lambda})^2\frac{1}{2}[\mathrm{Tr}(\mathcal{T}\mathcal{V}^\mu)\mathrm{Tr}(\mathcal{T}\mathcal{V}^\nu)]^2 \ , \\
\mathcal{L}_{11} &= \ell_{11}(\frac{v}{\Lambda})^2 \ g\epsilon^{\mu\nu\rho\lambda}\mathrm{Tr}(\mathcal{T}\mathcal{V}_\mu)\mathrm{Tr}(\mathcal{V}_\nu\mathbf{W}_{\rho\lambda}) \ ,
\end{aligned}
\tag{5}
$$

where $\mathbf{W}_{\mu\nu} = \partial_\mu\mathbf{W}_\nu - \partial_\nu\mathbf{W}_\mu + ig[\mathbf{W}_\mu, \mathbf{W}_\nu]$, $B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu$ and $\mathcal{V}_\mu \equiv (D_\mu U)U^\dagger$ , $\mathcal{T} \equiv U\tau_3 U^\dagger$ . The cutoff $\Lambda$ characterizes the scale for the effective

theory at which the new underlying dynamics sets in, presumably around $\Lambda = \min(4\pi v, M_{H,V})$. In such a normalization and without other symmetries in the underlying physics, one would expect the natural size for the couplings to be $\ell_i \sim \mathcal{O}(1)$.

### *2.4 Remarks on the anomalous couplings*

The operators in Eq. (5) modify the gauge boson interactions, leading to the so-called "anomalous couplings". Conventionally, the anomalous couplings of the gauge boson interactions are formulated by a Lorentz and electromagnetic gauge-invariant effective Lagrangian [22]

$$
\begin{aligned}
\frac{\mathcal{L}_V}{g_V} &= i g_1^V \left( W^\dagger_{\mu\nu} W^\mu V^\nu - W^\dagger_\mu V_\nu W^{\mu\nu} \right) + i\kappa_V W^\dagger_\mu W_\nu V^{\mu\nu} \\
&+ i \frac{\lambda_V}{m_W^2} W^\dagger_{\lambda\mu} W^\mu{}_\nu V^{\nu\lambda} - g_4^V W^\dagger_\mu W_\nu \left( \partial^\mu V^\nu + \partial^\nu V^\mu \right) \\
&+ g_5^V \epsilon^{\mu\nu\lambda\rho} \left( W^\dagger_\mu \partial_\lambda W_\nu - \partial_\lambda W^\dagger_\mu W_\nu \right) V_\rho + i\tilde{\kappa}_V W^\dagger_\mu W_\nu \tilde{V}^{\mu\nu} \\
&+ i \frac{\tilde{\lambda}_V}{m_W^2} W^\dagger_{\lambda\mu} W^\mu{}_\nu \tilde{V}^{\nu\lambda}.
\end{aligned}
\tag{6}
$$

In the SM at tree level, $g_1^V = \kappa_V = 1, \lambda_V = \tilde{\lambda}_V = \tilde{\kappa}_V = g_4^V = g_5^V = 0$. The deviation from the SM values can be expressed in terms of the coupling relations between this formalism and the electroweak chiral Lagrangian as

$$
\begin{aligned}
\Delta g_1^Z &= \frac{v^2}{\Lambda^2} \left( \frac{\ell_0}{c_{2w}} + \frac{e^2 \ell_1}{c_w^2 c_{2w}} + \frac{e^2 \ell_3}{c_w^2 s_w^2} \right), \qquad g_5^Z = \frac{v^2}{\Lambda^2} \frac{e^2 \ell_{11}}{c_w^2 s_w^2}, \\
\Delta\kappa_\gamma &= \frac{v^2}{\Lambda^2} \frac{e^2}{s_w^2} (-\ell_1 + \ell_2 + \ell_3 - \ell_8 + \ell_9), \\
\Delta\kappa_Z &= \frac{v^2}{\Lambda^2} \left( \frac{\ell_0}{c_{2w}} + \frac{2e^2 \ell_1}{c_{2w}} - \frac{e^2 \ell_2}{c_w^2} + \frac{e^2}{s_w^2} [\ell_3 - \ell_8 + \ell_9] \right),
\end{aligned}
\tag{7}
$$

where $s_w = \sin\theta_W,\ c_{2w} = \cos 2\theta_W$.

Note that these couplings are three-point interactions and always involve some pure gauge fields that may not be sensitive to the EWSB sector. This is especially true for those couplings involving photons, no matter it is a triple or quartic coupling. In contrast, the couplings $\ell_{4,5,6,7,10}$ are of particular interests: They are the *genuine* Goldstone boson interactions which characterize the underlying physics responsible for the EWSB.[13,23] Furthermore, those couplings are expected to be enhanced [13,24] over the gauge couplings due to the

new strong dynamics.[25] For instance,[26]

$$\begin{array}{ll} \text{heavy scalar } M_H = 2 \text{ TeV}: & \ell_4 \approx 0.0, \quad \ell_5 \approx 0.33; \\ \text{heavy vector } M_V = 2 \text{ TeV}: & \ell_4 \approx 0.38 \quad \ell_5 \approx -0.31. \end{array}$$

## 3 Current constraints on SEWS

Physics associated with vector boson pairs has been experimentally studied both at the Tevatron [27] and at LEP2.[28] The triple gauge-boson couplings can be measured through those processes. With good agreements with the SM expectation, the constraints from LEP2, CDF and D0 on the relevant anomalous couplings at a $2\sigma$ level are [28]

$$\left\{ \begin{array}{l} -0.07 < \Delta g_1^Z < 0.05 \implies -8.4 < \ell_3 < 4.5 \\ -0.11 < \Delta\kappa_\gamma < 0.23 \implies -16 < \ell_{2,9} < 35 \end{array} \right.$$

for $\Lambda = 2$ TeV. So far, the vector boson pairs produced are dominantly transversely polarized from the light fermion radiation, independent of the Higgs sector. Although the constraints will be further tightened up with more LEP2 data coming out, one may only expect to discover new physics signal through those channels when there is a resonance accessible or nearby.

Turning to the EWSB sector, besides the Higgs mass limit from the direct search,[1] precision electroweak data can be used to infer a Higgs mass limit $M_H < 107^{+67}_{-45}$ GeV,[29] which is only valid for a SM-like Higgs boson. Furthermore, from oblique corrections at the $Z$ pole, one can relate the $S, T, U$ parameters [5] to the anomalous couplings at tree level

$$S = -\frac{1}{\pi}(\frac{4\pi v}{\Lambda})^2 \ell_1, \qquad T = \frac{1}{2\pi e^2}(\frac{4\pi v}{\Lambda})^2 \ell_0, \qquad U = -\frac{1}{\pi}(\frac{4\pi v}{\Lambda})^2 \ell_8, \tag{8}$$

and translate the limits on $S, T, U$ [29] to

$$\left\{ \begin{array}{l} S = -0.27 \pm 0.12 \implies 0.04 < \ell_1 < 0.67 \\ T = 0.00 \pm 0.15 \implies -0.08 < \ell_0 < 0.08 \\ U = 0.19 \pm 0.21 \implies -0.80 < \ell_8 < 0.30 \end{array} \right.$$

for $\Lambda = 2$ TeV and at $2\sigma$. We see that the constraint from $T$ on the custodial symmetry breaking effect is very strong.[b] Loop corrections to the $T$ parameter through other couplings have been evaluated.[31] With the new $T$ value above, the limits on the couplings of interest are

$$\begin{array}{lll} -11 < \ell_4 < 11, & -28 < \ell_5 < 28, & -1.4 < \ell_6 < 1.4, \\ -11 < \ell_7 < 11, & -1.5 < \ell_{10} < 1.5. & \end{array}$$

[b] A new analysis appeared more recently in this topic.[30]

This calculation is in principle similar to that by considering the $Z$ partial widths.[32]

It is interesting to note that the couplings $\ell_{6,7,10}$ violate the custodial SU(2) symmetry, yet the current constraints on them are not very strong. It would be of great theoretical significance if one observed the custodial SU(2) symmetry breaking in SEWS. Currently, low energy constraints on heavy resonances are still weak with one exception for a vector resonance like $\rho_T$ which may mix with $\gamma, Z$. However, if there is also a nearly degenerate axial vector $A_1$, then the constraint on $\rho_T$ can be avoided.[19]

## 4 Quartic gauge-boson couplings at linear colliders

Physics potential at $e^+e^-$ colliders has been nicely presented in recent review articles [33] and in many talks in this workshop. To explore the physics of the electroweak symmetry breaking sector, the *most direct* way is to study the four-gauge boson couplings to probe the genuine Goldstone boson interactions, no matter there is a light Higgs boson or not. The relevant processes and the corresponding couplings of our interests are

$$\begin{aligned} e^+e^- &\to W^+W^-Z, \quad (\ell_{4,5,6,7}) && (9)\\ e^+e^- &\to ZZZ. \qquad (\ell_{4,5,6,7,10}) && (10)\end{aligned}$$

which are more useful at lower energies near the three-vector boson threshold. At higher energies, the fusion processes $W_L^* W_L^* \to W_L W_L$ become dominant

$$\begin{aligned} e^+e^- &\to \bar\nu\nu W^+W^-, \quad (\ell_{4,5}, H, \rho_T) && (11)\\ e^+e^- &\to \bar\nu\nu ZZ, \qquad (\ell_{4,5}, \ell_{6,7}, H) && (12)\\ e^+e^- &\to e^\pm\nu W^\mp Z, \qquad (\ell_{4,5}, \ell_{6,7}, \rho_T) && (13)\\ e^+e^- &\to e^+e^- ZZ. \qquad (\ell_{4,5} + 2\ell_{6,7,10}) && (14)\end{aligned}$$

With the linear collider running at different beam modes, certain other processes can be complementary for the study

$$\begin{aligned} e^-e^- &\to \nu\nu W^-W^-, \qquad (\ell_{4,5}, I=2) && (15)\\ \gamma\gamma &\to W^+W^-ZZ,\ W^+W^-W^+W^-, && (16)\\ e\gamma &\to eW^+W^-,\ eZ\gamma,\ eZZ. \quad (\omega_T, A_1) && (17)\end{aligned}$$

If the top quark plays a role in the EWSB, then the direct probe to the physics associated with this sector may be via the subprocess $W_L W_L \to t\bar t$ for

$$e^+e^- \to \bar\nu\nu W^* W^* \to \bar\nu\nu \bar t t. \qquad (H, \rho_T ...) \qquad (18)$$

We now summarize the current results for studies of the above processes.

*4.1 Triple gauge-boson production*

Triple gauge-boson production at $e^+e^-$ linear colliders will be an ideal process to study the quartic gauge boson couplings.[34,35,26] As outlined in Eqs. (9)-(10), the processes $e^+e^- \to W^+W^-Z$ and $ZZZ$, involving $WWZZ, ZZZZ$ couplings, can be sensitive to new physics. In Fig. (1), we present the sensitivity contours [26] at 90% C.L. for these two processes at different energies. We see that with $\sqrt{s} = 500$ GeV and an integrated luminosity of 50 $\text{fb}^{-1}$, the magnitude of the couplings can be probed to about $4-10$. At $\sqrt{s} = 1.6$ TeV with 200 $\text{fb}^{-1}$, the sensitivity to the couplings can be reached to a level of $0.3-0.6$, reaching theoretically quite interesting region. Beam polarizations of 90% for $e^-$ and 65% for $e^+$ helped to reduce the SM background.

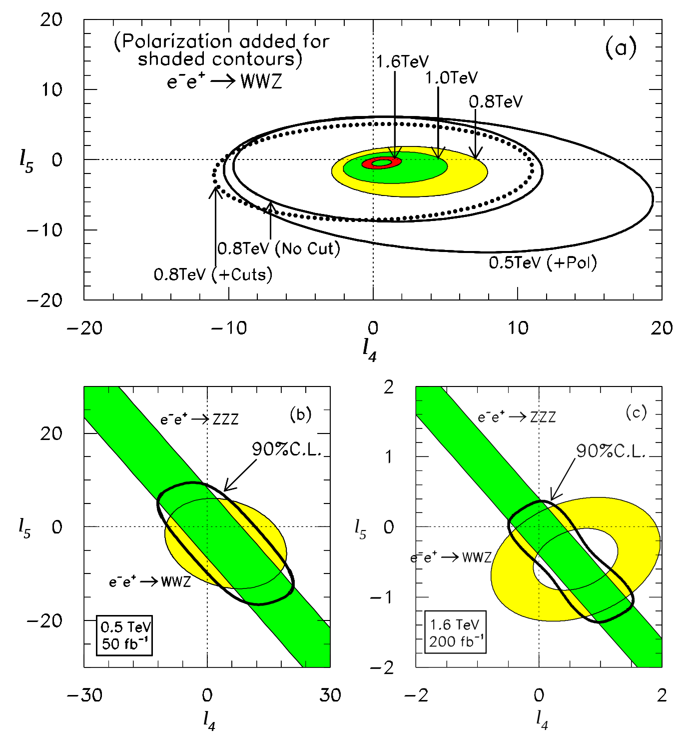


Figure 1: $1\sigma$ contours in $\ell_4 - \ell_5$ plane via $WWZ, ZZZ$ final states. We take $\Lambda = 2$ TeV. The thick solid lines are for the two-channel combined 90% C.L. bounds.

### *4.2 WW fusion processes*

Cross sections for the processes of Eqs. (9)-(10) are suppressed at higher energies, typically like $s^{-1}$ well above the threshold. On the other hand, the gauge boson fusion processes of Eqs. (11)-(14), originated by the subprocesses $W^+W^- \to W^+W^-, ZZ; W^\pm Z \to W^\pm Z; ZZ \to ZZ$ and $W^-W^- \to W^-W^-$, become more significant, with an energy dependence typically like $\ln(s/M_W^2)$ due to the collinear weak boson radiation off the electron beams in the initial state. In particular, if there are new resonances in the SEWS sector that are accessible at the collider, then the signal should be more substantial. The backgrounds to the gauge-boson fusion and the signal isolation technique have been extensively studied first at hadron colliders [36,18] and later at the linear colliders.[37] The SM backgrounds and SEWS signals are calculated for a 1.5 TeV linear collider for a scalar (Higgs-like), a vector ($\rho_T$-like) of a mass one TeV, and the low-energy theorem amplitude (LET, non-resonance), which are shown in Fig. 2. The signal is clearly observable above the backgrounds after judicious cuts. It is important to note that by examining the individual $W^+W^-$ or $ZZ$ final state, one may be able to deduce the structure of the underlying dynamics. For example, the ratio $\sigma(W^+W^- \to W^+W^-)/\sigma(W^+W^- \to ZZ)$ is about two for a Higgs-like scalar, is much larger than one for a vector resonance, and is even smaller than one for the LET.[37]

For non-resonant scenario beyond the LET, the fusion processes are also studied [38] for the couplings $\ell_{4,5,6,7,10}$, as shown in Fig. 3. The characteristic range of the probe for a 1.6 TeV collider with 200 fb$^{-1}$ luminosity is about 0.06, which goes well below unity. This sensitivity would be of great theoretical interest.

### *4.3* $W_L W_L \to t\bar{t}$

The possibly significant role of the top quark played in the EWSB sector motivates the study of the process Eq. (18). The enhancement of the cross section due to the contribution from lower-lying resonances can be quite substantial and is shown in Fig. 4 versus the $e^+e^-$ c.m. energy.[39] We view the SM result (with $m_H = 100$ GeV) as the background. The signal rates are evaluated with the approximation of the Goldstone-boson equivalence theorem, and are labeled by $M_H = 1$ TeV (a heavy scalar), $M_V = 1$ TeV (a heavy vector) and LET (leading order non-resonance). Another study on this channel was reported during the workshop [40] and good signals in the $M_{t\bar{t}}$ distributions for a scalar/vector resonances could be found. The statistical significance for the signal channels considered can be well above 10.

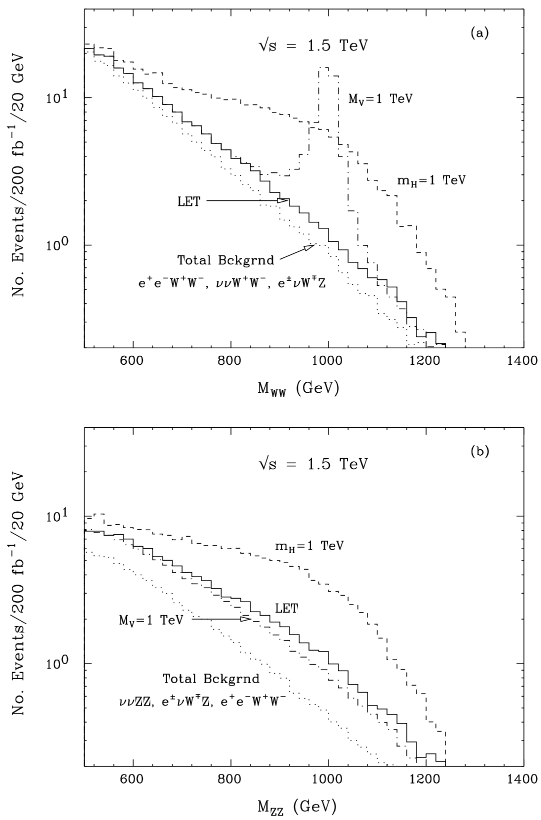


Figure 2: Scalar and vector resonant production via (a) $W_L^+W_L^- \rightarrow W_L^+W_L^-$ and (b) $W_L^+W_L^- \rightarrow ZZ$ at a 1.5 TeV $e^+e^-$ linear collider. The summed SM background (dotted) and a LET amplitude (solid) are also presented.

## 5 Concluding remarks

A few remarks are in order. First of all, regarding the machines. We see that a TeV $e^+e^-$ linear collider should have great potential for probing EWSB physics. On the other hand, $e^-e^-$, $e^-\gamma$ and $\gamma\gamma$ colliders can be all complementary. As indicated in Eqs. (15)-(17), the $e^-e^-$ collision is unique in probing the weak isospin $I = 2$ $W^-W^-$ scattering;[41] $e\gamma$ is unique in producing the weak isosinglet states such as $\omega_T$ via $\gamma Z$ fusion;[21] and a $\gamma\gamma$ collider has similar

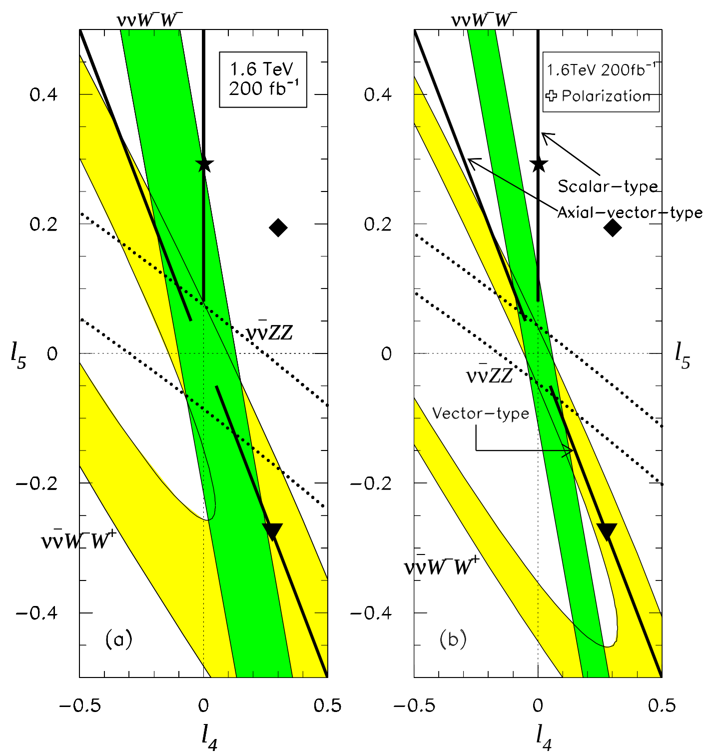


Figure 3: $1\sigma$ contours in $\ell_4 - \ell_5$ plane via $WW$ fusion processes with $W^+W^-$, $ZZ$ and $W^-W^-$ final states. We take $\Lambda = 2$ TeV.

potential to the $e^+e^-$ collider.[42] Secondly, about the detectors. To effectively distinguish a $W$ from a $Z$ in the hadronic modes, one would need an adequate hadronic mass resolution for the calorimeter. The detector coverage in the forward region should be at least of the order of 10 degrees in order to tag/veto the forward leptons for the fusion processes. Not mentioned in the presentation is the situation where a vector resonance mixes with $\gamma/Z$ and the mass is close or below the energy threshold. In this case, the signal would be particularly strong[43] and one can study its properties to a great detail.

Finally, we provide a "naive" comparison between a 1.5 TeV linear collider [37,26,38] with a luminosity of 200 fb$^{-1}$ and the LHC [18,44] at 14 TeV with 100 fb$^{-1}$. We see from Table 1 that they have comparable reach for most of cases

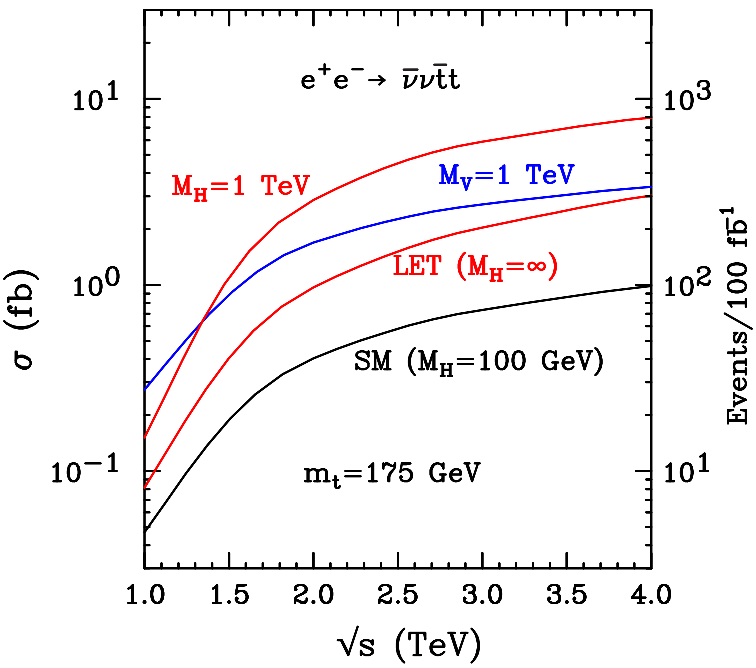


Figure 4: Production cross section versus the c.m. energy for $e^+e^- \to \bar{\nu}\nu\bar{t}t$. The SM rate as the background is that of $m_H = 100$ GeV; the signal rates are labeled by $M_H = 1$ TeV (a heavy scalar), $M_V = 1$ TeV (a heavy vector) and LET (leading order non-resonance).

under discussion. Due to the advantage of a cleaner experimental environment at the linear collider, the $t\bar{t}$ channel seems to be more accessible there than at the LHC.

Table 1: Comparison for the sensitivity for SEWS physics between a 1.5 TeV $e^+e^-$ collider and the LHC. The entries under scalar, vector and $t\bar{t}$ are for the estimated statistical significance; those under $\ell_{4,5}$ are the values that can be probed for the anomalous couplings.
.

| | scalar(1 TeV) | vector(1 TeV) | $\ell_{4,5}$ | $t\bar{t}$ |
|---|---|---|---|---|
| LC(1.5 TeV) | 10 | 16 | $-0.1-0.1$ | 10 |
| LHC(14 TeV) | 8 | 15 | $-0.5-0.9$ | not established |

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